Conductance Fluctuations of Mesoscopic Spin Glasses

P. G. N. de Vegvar, L. P. Lévy, and T. A. Fulton
AT&T Bell Laboratories, Murray Hill, New Jersey 07974
(Received 3 December 1990)

Magnetoresistance measurements on 1000-ppm *Cu*Mn mesoscopic spin-glass samples reveal phase-coherent electron propagation on 0.4-μm length scales observed as conductance fluctuations and Aharonov-Bohm resistance oscillations. Exploiting the broken time-reversal symmetry of the spin-glass state, a unique magnetofingerprint of the specific frozen spin configuration is extracted from a pair of resistance measurements. These "spin-glass fingerprints" are manifestly different after several thermal histories, but are surprisingly robust upon field cycling to over 9 times the typical exchange field.

PACS numbers: 72.15.Rn, 72.15.Qm, 75.50.Lk

In disordered metals at sufficiently low temperatures, electrons diffuse over large distances (l_ϕ) before losing their phase coherence. The resulting quantum interferences between diffusive paths depend sensitively on the specific locations of impurities. This is manifested in the magnetoconductance of small samples[1] which exhibits aperiodic fluctuations as the electronic phases change each time a flux quantum $\Phi_0 = h/e$ is threaded through a phase-coherent region. Such magnetofingerprints of impurity configurations have been proposed as a microscopic probe of the complex *frozen* magnetic order in a metallic spin glass.[2] In principle, different spin configurations may be distinguished by the differences in their magnetoconductances yielding information on the accessible spin-glass states and some measure of their overlap. Since the effective sample size can be easily changed by altering the voltage probe spacing, any characteristic length scale of the spin-glass phase, which may be a correlation length or a size for droplet excitations,[3] could also be isolated if a measurable phase-coherent fingerprint were to be demonstrated. This presupposes, however, that local fields are sufficiently large to freeze-out spin-flip dephasing, a point questioned by recent experimental studies.[4]

These difficulties are overcome in the experiments described in this Letter, where we report observations of phase-coherent transport in 1000-ppm mesoscopic *Cu*Mn structures. Exploiting the broken time-reversal symmetry of the spin-glass phase, we also show how the spin-dependent contributions to the conductance fluctuations can be extracted to obtain a "fingerprint" of a spin-glass configuration. We then examine how this fingerprint changes with temperature and field. As the system is warmed through T_g the spin fingerprint disappears, and upon refreezing a new one arises as the system settles into a new state. Yet, most surprisingly, these spin fingerprints are remarkably unchanged on field cycling to many times the exchange field $h_g = k_B T_g/\mu$.

To obtain a long phase breaking length l_ϕ, a dilute alloy [1000-atomic-ppm *Cu*Mn, with inter-Mn spacing $d \approx 23$ Å, bulk $T_g \approx 1.6$ K, and bulk residual resistivity $\rho \approx 4.3$ $\mu\Omega$ cm/at.% (Ref. 5)] was selected for its large diffusion constant D, weak spin-orbit scattering, and small out-of-phase susceptibility $\chi''(\omega)$ reflecting a more rigidly frozen spin system. *Cu*Mn films, wires, and rings (with linewidths ≤ 850 Å, see Fig. 1 inset) were deposited to depths of 420–500 Å using two different methods: a low-pressure argon-ion sputtering technique and a flash-evaporation method.[6] The data were qualitatively similar for all samples. The longest elastic mean free path[7] $\lambda_e \approx 310$ Å was obtained for sputtered samples whose physical properties are now described.

The temperature dependence of the resistance R in zero applied magnetic field for a sputtered wire and a codeposited film are shown in Fig. 1. The decrease of R with increasing T above the bulk T_g is nearly logarithmic and is most likely due to the Kondo effect.[8] For $2 < T < 8$ K, we find that $\Delta\rho(T) = -A(T)\ln(T)$, with $A(T)$ varying slowly with T: $A(\text{film}) = 3.0$ nΩ cm,

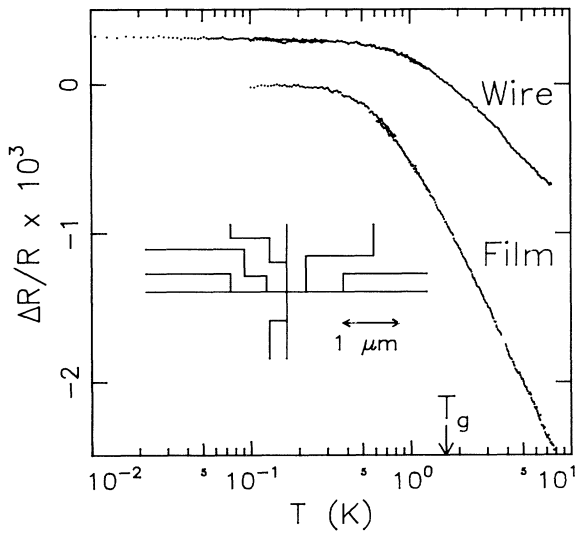

FIG. 1. Temperature dependence of resistance for a 420-Å-thick by 2-μm-wide 1000-ppm *Cu*Mn film and a codeposited 850–900-Å-wide by 2-μm-long wire in zero applied field. The bulk T_g is shown for comparison. Inset: A scaled schematic of one wire sample.

A(wire) = 1.2 nΩ cm. Comparing to measured values of A = 4.0 nΩ cm on bulk samples,[5] these $R(T)$ are consistent with Kondo behavior suppressed by interactions between impurities.[9] Below T_g, however, the resistance of both sputtered and flash-evaporated films and wires gradually saturates. This contrasts markedly with the decrease in resistance observed in bulk samples as T drops below T_g.[5] The bulk behavior is usually attributed to the freeze-out of single spin-flip scattering as the moment fully polarizes in its local field. Within this picture, the saturation of the film and wire resistances below $0.6T_g$ implies a total polarization of the Mn impurities and an unexpectedly large suppression of spin-flip scattering relative to the bulk. Although we presently lack a complete understanding of this suppression, the observability of the mesoscopic effects discussed below hinges on reducing spin-flip dephasing for $T < T_g$ from its estimated bulk value.

The large negative magnetoresistance shown in Fig. 2 for a film can be used to infer that these samples' magnetotransport properties are similar to bulk spin glasses. The magnitude of the magnetoresistance is comparable to that observed in bulk samples after scaling by the ratio of magnetic to nonmagnetic scattering[5] and is too large ($8e^2/h$) to be attributed to weak localization.[10] The near saturation at high fields suggests that the negative magnetoresistance is due to a change of the electron elastic-scattering rate as the Mn ions polarize along the applied field.[11] At low temperatures, the inflection point of the magnetoresistance gives a characteristic field of order 9–11 kG, comparable to the typical exchange field $h_g \approx k_B T_g/\mu \approx 8$ kG expected for a spin-glass order.

Figure 2 also displays the magnetoresistance of a 2-μm-long wire segment. The observed quantum fluctuations (a magnetofingerprint) arise from two sources: (a) The magnetic field induces a phase shift in the electronic wave functions, giving rise to fluctuations on a field scale $h_\phi = \Phi_0/wl_\phi$, where w is the wire width. (b) A spin configuration has a finite susceptibility and is distorted by the external field: This induces conductance fluctuations on a field scale h_s, which is determined below to be significantly larger than h_ϕ. The magnetofingerprints are therefore dominated by the *frozen* initial configuration, and not by the spin-glass stiffness. At 23 mK the conductance fluctuations have a peak-to-peak amplitude $\Delta G = 0.20e^2/h$ at high fields. Their magnitude does not change up to about 1.0 K and decreases gradually thereafter. The fluctuations possess a correlation coefficient $C = 0.94$ upon cycling the magnetic field to 70 kG: The

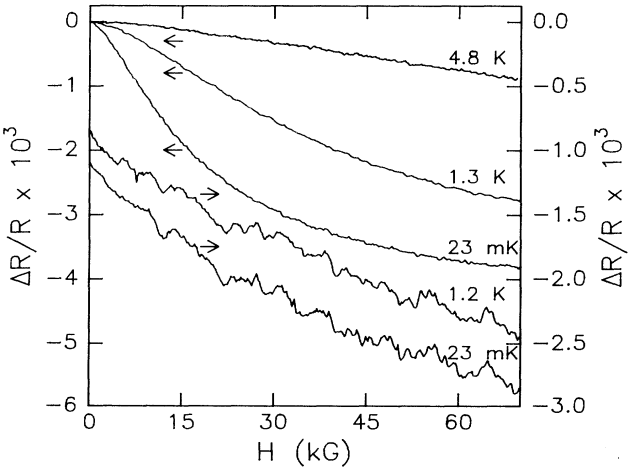

FIG. 2. The top three curves (left vertical scale) display the magnetoresistance of the same film as Fig. 1 at different temperatures. The lower two traces (right vertical scale) show the magnetoresistance of the 2-μm-long wire.

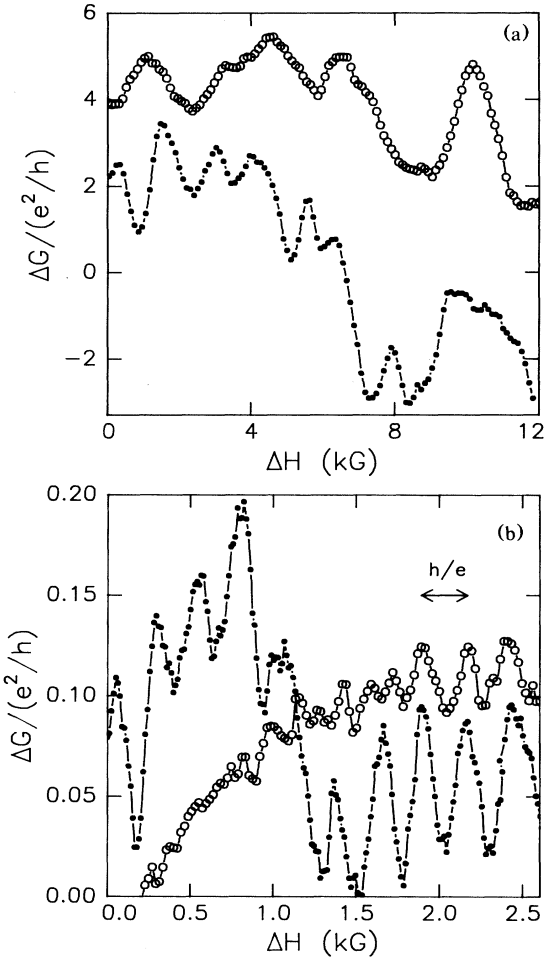

FIG. 3. (a) Conductance fluctuations in a wire segment with a voltage probe spacing of 0.35 μm at 23 mK. The open (solid) points refer to data acquired over a −6–6-kG (58–70-kG) field range. (b) Aharonov-Bohm conductance oscillations as measured in a square ring 0.4 μm on a side at 23 mK. The open (solid) data were collected from 1 to 3.5 kG (50 to 52.5 kG).

spin configurations are therefore quite robust upon field cycling to $9h_g$. The conductance fluctuations, which are expected in the fully polarized limit ($h \gg h_g$),[12] persist into the spin-glass regime at fields below h_g, where the larger spin-flip dephasing rate could destroy the quantum effects. This is apparent from Fig. 3(a), where the magnetoresistance of a 0.35-μm wire section measured between 58–70 and ± 6 kG is plotted. These data are highly repeatable; two successive low-field scans have $C=0.992$. The most striking demonstration of phase-coherent transport is given by the Aharonov-Bohm resistance oscillations measured in a square-shaped ring 0.4 μm on a side, displayed in Fig. 3(b). The oscillations decrease in amplitude as the field is dropped from 50–52 down to 1–3.5 kG, but do not disappear. Since the conductance oscillates with a period Φ_0, the field $h_\phi \approx 260$ G controls the conductance fluctuations and is therefore significantly smaller than h_s for the sample studied.

The dephasing length l_ϕ may by estimated in two ways. From the decorrelation field $h_\phi = \Phi_0/wl_\phi$ measured on a $L=0.35$ μm segment at 22 mK, we find a monotonic increase of l_ϕ from 0.37 μm below 10 kG to 0.96 μm at 60 kG. Because the Aharonov-Bohm oscillation peak-to-peak amplitude ΔG decreases exponentially with the ring perimeter[12] L as $\Delta G \approx 0.4(e^2/h)\exp(-L/2l_\phi)$, we find from Fig. 3(b) that l_ϕ ranges from 0.32 μm at 1–4 kG to 1.2 μm at 50 kG, in agreement with the first estimate.

Since both magnetic and nonmagnetic scattering contribute to the conductance fluctuations, information about the specific spin configuration in the sample can be obtained only by separating the spin-dependent and spin-independent contributions to the magnetoresistance. This is particularly important here since the fraction of scattering (potential and spin dependent) that occurs from the Mn spins is only 15% due to scattering from grain boundaries and surfaces. To this end, we recall the Onsager-Büttiker relations[13]

$$R_{i,v}(H,M) = R_{v,i}(-H,-M)\,, \tag{1}$$

which are valid only for a system invariant under time reversal. Here i,v refers to a set of current and voltage leads used to perform the measurement and M is the magnetization. Because a frozen magnetic state breaks time-reversal symmetry, the Onsager-Büttiker antisymmetrized resistance $R_a = \frac{1}{2}[R_{i,v}(H) - R_{v,i}(-H)]$ no longer vanishes and represents the contribution from the magnetic order. Actually, this magnetic contribution can be further decomposed into one component $R_{a,s}$ symmetric in H and another one $R_{a,a}$ antisymmetric in the field. Here we will focus on the former, defined as

$$R_{a,s}(H) = \tfrac{1}{4}[R_{i,v}(H) - R_{v,i}(-H) + R_{i,v}(-H) - R_{v,i}(H)]\,. \tag{2}$$

This quantity is depicted in Fig. 4 for a 0.65-μm-long segment as a conductance fluctuation, where the data

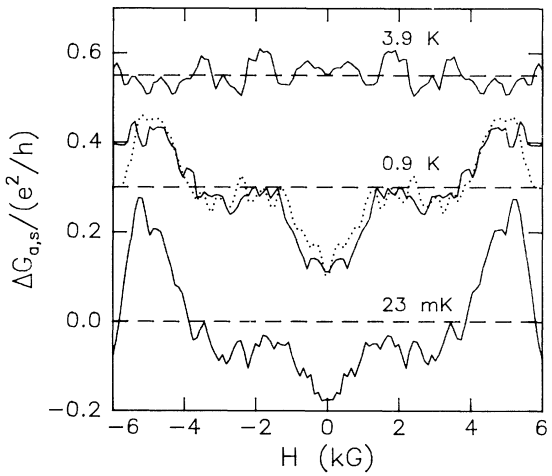

FIG. 4. The $\Delta G_{a,s}$ component (see text) of the magnetoconductance fluctuations as extracted from measurements on a segment with a 0.65-μm voltage probe spacing. The sample was cooled from 4 K to 12 mK in 7 kG and subsequently warmed to the indicated temperatures. The curves are offset vertically for clarity. The two traces at 0.9 K indicate the experimental reproducibility and have $C=0.90$.

were obtained on warming from 23 mK. The 23-mK trace is very similar to the 0.9-K ones with an average $C=0.89$. This implies that there is little change in the magnetic state on this length scale below 0.9 K. As the sample is warmed further through the bulk T_g to 2.2 K (not shown), $\Delta G_{a,s}$ is only partially correlated with the 0.9-K curves, having $C=0.49$. By $T=3.9$ K, this component of the conductance contrasts vividly with the lower-temperature data: It has dropped into the noise floor of the measurement. This behavior is expected if the spins start to lose their rigidity at T_g. On the other hand, the Onsager-Büttiker symmetric components of the conductance fluctuations did not display such a sharp decrease in amplitude and were highly correlated across T_g with a mean $C=0.85$. Similarly, the fluctuations in $R_{a,s}$ and $R_{s,s}$ have different temperature dependences: The rms fluctuations of the former (latter) drop by a factor of 0.42 (0.88) from 23 mK to 2.2 K. The association of this "melting" behavior observed in a wire with the bulk T_g also indicates the Mn impurities in the wire are not appreciably oxidized.

We now determine the sensitivity of the conductance to a specific spin configuration; i.e., how many spins must reorient before $\Delta G_{a,s} = e^2/h$? This sensitivity arises in the following way. When $L < l_\phi$, the resistance measured by two voltage leads may be viewed as a combination of transmission probabilities for an electron to be scattered from one lead into another.[14] Each transmission amplitude between an incoming (i) and an outgoing (o) channel can be thought of as a sum over amplitudes corresponding to Feynman paths. Because of the

diffusive motion, such paths are random walks between i and o channels consisting of segments of tubes having radius k_F^{-1}, the inverse Fermi wave vector, and of length λ_e. Any given path visits a large number of impurities, occupying a volume $v \sim (4\pi\lambda_e/k_F^2)(L/\lambda_e)^2$. If the scattering potential of a single impurity changes, it can induce a conductance change $\Delta G \approx (e^2/h)\sqrt{f}$, where $f = v/V$ is the fraction of Feynman paths visiting the impurity and V the sample volume.[15] For a wire of length $L = l_\phi \approx 0.35$ μm, the number of Mn impurities contained in this Feynman path volume $v \sim 2.7 \times 10^5$ Å^3 is $N = v/d^3 = 20$. Taking into account the fraction of scattering due to the Mn, we find that $\sim 28\%$ of the 1.1×10^5 spins in V must reorient to produce $\Delta G = e^2/h$, provided each reoriented spin produces a relative phase shift of order π. For $L = l_\phi$ we can detect $\Delta G = 0.07 e^2/h$. For lengths $L > l_\phi$, ΔG will be reduced by a factor $(L/l_\phi)^{-3/2}$ due to the random addition of each l_ϕ segment.

This technique may then be applied to measure the extent to which the spin configuration changes with T or H. To investigate the former we measured the low-field magnetoresistances of a $L = 2$ μm wire at $T/T_g \approx 0.4$ separated by an anneal to $T/T_g = 2.2$ in $H = 0$ kG and found an rms $(L/l_\phi)^{3/2} \Delta G_{a,s} \approx 0.33 e^2/h$.[16] This value implies a large fraction of the spins in an l_ϕ long volume reoriented due to the anneal. Such a reorientation is reflected in the two $\Delta G_{a,s}(H)$ spin fingerprints, which are typically only weakly correlated across the anneal: $C = 0.34$–0.68. This behavior stands in stark contrast to cycling H starting from a zero-field-cooled configuration to 75 kG $\approx 9 h_g$ and back at fixed $T = 0.8$ K. Then the low-field fingerprints had $C = 0.90$, implying a *nearly reversible* distortion of the spin configuration as the field was cycled; i.e., the magnetic state of the system remained in nearly the same energy valley even though it had been polarized by H.[17]

To conclude, we have demonstrated that phase-coherent transport is experimentally observable in a spin glass and that information about the spin configuration on short length scales can be obtained from a set of magnetoresistance measurements. We have illustrated the method by showing that warming through T_g changes the magnetic state far more effectively than field cycling. This powerful technique can in principle provide quantitative information on the overlaps between spin-glass configurations when the theoretical connection between magnetoconductance traces and configuration overlaps becomes elucidated.

We have benefited from conservations with S. Hershfield and B. L. Al'tshuler. We also gratefully acknowledge L. N. Dunkleberger, G. Espinosa, G. Kammlott, R. Miller, and P. Trevor for their contributions to the processing and C. A. Murray for the UHV system.

[1]C. P. Umbach, S. Washburn, R. B. Laibowitz, and R. A. Webb, Phys. Rev. B **30**, 4048 (1984); P. A. Lee and A. D. Stone, Phys. Rev. Lett. **55**, 1622 (1985); B. L. Al'tshuler, Pis'ma Zh. Eksp. Teor. Fiz. **41**, 530 (1985) [JETP Lett. **41**, 648 (1985)].

[2]B. L. Al'tshuler and B. Z. Spivak, Pis'ma Zh. Eksp. Teor. Fiz. **42**, 363 (1985) [JETP Lett. **42**, 447 (1985)]; S. Feng, A. J. Bray, P. Lee, and M. A. Moore, Phys. Rev. B **36**, 5624 (1987); N. E. Israeloff, M. B. Weissman, G. J. Nieuwenhuys, and J. Kosiorowska, Phys. Rev. Lett. **63**, 794 (1989).

[3]L. P. Lévy and A. T. Ogielski, Phys. Rev. Lett. **57**, 3288 (1986); G. G. Kenning *et al.*, Phys. Rev. B **42**, 2393 (1990); H. Vloeberghs, J. Vranken, C. Van Haesendonck, and Y. Bruynseraede, Europhys. Lett. **12**, 557 (1990); A. Yu. Zyuzin and B. Z. Spivak, Pis'ma Zh. Eksp. Teor. Fiz. **43**, 185 (1986) [JETP Lett. **43**, 234 (1986)]; A. Jagannathan, E. Abrahams, and M. Stephens, Phys. Rev. B **37**, 436 (1988); D. Fisher and D. Huse, Phys. Rev. B **38**, 386 (1988).

[4]C. Van Hasendonck, H. Vloeberghs, and Y. Bruynseraede, in *Nanostructure Physics and Fabrication,* edited by M. Reed and W. Kirk (Academic, New York, 1989); A. Benoit and V. Chandrasekhar (private communication).

[5]J. Bass, in *Metals: Electronic Transport Phenomena,* edited by J. Bass and K. H. Fischer, Landolt-Börnstein New Series, Group 3, Vol. 15, Pt. a (Springer, Berlin, 1983), p. 166; O. Laborde and P. Radhakrishna, J. Phys. F **3**, 1731 (1973); P. J. Ford and J. A. Mydosh, Phys. Rev. B **14**, 2057 (1976); A. K. Nigam and A. K. Majumdar, Phys. Rev. B **27**, 495 (1983).

[6]Bi-layer Cr-polymethylmethacrylate resist and e-beam lithography were used: G. J. Dolan and J. Dunsmuir, Physica (Amsterdam) **152B**, 7 (1980). The material was deposited at a rate of 0.2 Å/sec in a UHV sputtering system [$p \approx (1$–$2) \times 10^{-9}$ Torr] to avoid Mn oxidation.

[7]The corresponding 4.2-K resistivity was 2.1 $\mu\Omega$ cm.

[8]R. P. Peters, G. Bergmann, and R. M. Mueller, Phys. Rev. Lett. **58**, 1964 (1987); P. Monod, Phys. Rev. Lett. **19**, 1113 (1967).

[9]A is 2 orders of magnitude too large to be attributed to weak localization or interaction effects.

[10]Weak localization is actually destroyed in vector spin glasses: Wei Wei and G. Bergmann, Phys. Rev. B **38**, 11751 (1988); L. P. Lévy and P. G. N. de Vegvar, Bull. Am. Phys. Soc. **34**, 415 (1989).

[11]M. T. Beal-Monod and R. A. Weiner, Phys. Rev. **170**, 552 (1968); K. H. Fisher, Z. Phys. B **42**, 27 (1983).

[12]S. Washburn and R. A. Webb, Adv. Phys. **35**, 375 (1986).

[13]M. Büttiker, IBM J. Res. Dev. **32**, 317 (1988), and references contained therein.

[14]R. Landauer, IBM J. Res. Dev. **32**, 306 (1988).

[15]P. A. Lee, in *Proceedings of the Eighteenth International Conference on Low Temperature Physics, Kyoto, 1987* [Jpn. J. Appl. Phys. **26**, Suppl. 26-3, 1934 (1987)]; B. L. Al'tshuler, *ibid.*, p. 1938.

[16]We also carried out a high-field (70-kG) control measurement to demonstrate that this change was not due to impurity diffusion.

[17]K. Binder and A. P. Young, Rev. Mod. Phys. **58**, 801 (1986); G. Parisi, M. Mezard, and M. Virasoro, *Spin-Glasses and Beyond* (World Scientific, Singapore, 1987).